

Strong CP Problem, Neutrino Masses and the 750 GeV Diphoton Resonance

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We present an $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_L \times U(1)_R \rightarrow SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ left-right symmetric model with a discrete parity symmetry to realize a universal seesaw scenario. The model can simultaneously solve the strong CP problem without resorting to the unobserved axion and explain the 750 GeV diphoton resonance reported recently by the ATLAS and CMS collaborations at the LHC. Owing to large suppressions in the two-loop induced Dirac mass terms, the Majorana mass matrices of left- and right-handed neutrinos naturally share the same structure. That allows us to quantitatively study the neutrinoless double beta decay induced by the right-handed currents.

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I. INTRODUCTION

The discovery of the 125 GeV Higgs boson at the LHC [1, 2] is a landmark which indicates the underlying reason of the spontaneous electroweak symmetry breaking and the mass generation mechanism. One question often raised is, in addition to the Higgs boson, whether or not there are other fundamental particles or mechanisms also play roles in the electroweak symmetry breaking. Especially, the smallness of neutrino masses is a puzzle which drives us to go beyond the standard model (SM). The seesaw mechanism [3–6] provides a natural way to understand the tiny size of the observed neutrino masses. Generalized seesaw mechanism (called universal seesaw) [7–10] were introduced to generate quark and charged lepton masses under a “seesaw” way, with an extension of a new set of vector-like singlet fermions (quarks and leptons). The universal seesaw can be elegantly achieved in some $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ left-right symmetric models [11] where the spontaneous breaking of left-right and electroweak symmetries are driven by two $[SU(2)]$ -doublet Higgs scalars, while the ordinary charged fermions from $SU(2)$ doublets obtain their masses by integrating out additional $[SU(2)]$ -singlet charged fermions. In such left-right symmetric models for the universal seesaw, one can further solve the strong CP problem without resorting to the unobserved axion if the discrete parity symmetry is imposed [12, 13].

In this work, we shall consider an $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_L \times U(1)_R$ left-right symmetric model where the $[SU(2)]$ -singlet fermions for the universal seesaw can naturally acquire their masses after a spontaneous symmetry breaking $U(1)_L \times U(1)_R \rightarrow U(1)_{B-L}$. For this symmetry breaking, we introduce some $[SU(2)]$ -singlet Higgs scalars which provide a natural solution to the 750 GeV diphoton resonance observed recently at

the LHC Run-2 with a center-of-mass energy $\sqrt{s} = 13$ TeV [14, 15]. It has drawn lots of interests in the field [16–87]. We further add two $[SU(2)]$ -triplet Higgs scalars to generate the Majorana mass matrices of the left- and right-handed neutrinos. The Dirac mass term between the left- and right-handed neutrinos can be highly suppressed since it is produced only by the charged gauge bosons and fermions at two-loop level. The right-handed neutrino mass matrix is proportional to the left-handed neutrino mass matrix so that the neutrinoless double beta decay from the right-handed currents can be quantitatively studied [88]. Our model can also accommodate the parity symmetry for solving the strong CP problem without an axion.

The paper is organized as follows: In Sec. II and III we present our model and the pattern of the symmetry breaking. In Sec. IV and V, we address on the charged fermion mass and neutrino mass generation and how to solve the strong CP problem. In Sec. VI we interpret the 750 GeV diphoton resonance as the singlet scalar, the mass generator of the vectorquark in our model. Finally, we conclude in Sec. VII.

II. THE MODEL

In the fermion sector we have the following $SU(2)$ singlets,

$$U_L(3, 1, 1, -\frac{4}{3}, 0) \xleftrightarrow{P} U_R(3, 1, 1, 0, -\frac{4}{3}), \quad (1a)$$

$$D_L(3, 1, 1, +\frac{2}{3}, 0) \xleftrightarrow{P} D_R(3, 1, 1, 0, +\frac{2}{3}), \quad (1b)$$

$$E_L(1, 1, 1, +2, 0) \xleftrightarrow{P} E_R(1, 1, 1, 0, +2), \quad (1c)$$

in addition to the $SU(2)$ doublets as below,

$$\begin{array}{c} q_L(3, 2, 1, +\frac{1}{3}, 0) = \begin{bmatrix} u_L \\ d_L \end{bmatrix} \\ \uparrow P \downarrow \\ q_R(3, 1, 2, 0, +\frac{1}{3}) = \begin{bmatrix} u_R \\ d_R \end{bmatrix}, \end{array} \quad (2a)$$

$$\begin{array}{c} l_L(1, 2, 1, -1, 0) = \begin{bmatrix} \nu_L \\ e_L \end{bmatrix} \\ \uparrow P \downarrow \\ l_R(1, 1, 2, 0, -1) = \begin{bmatrix} N_R \\ e_R \end{bmatrix}. \end{array} \quad (2b)$$

Here and thereafter the brackets following the fields describe the transformations under the $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_L \times U(1)_R$ gauge groups. As for the Higgs scalars, they include three $SU(2)$ singlets,

$$\sigma_U(1, 1, 1, +\frac{4}{3}, -\frac{4}{3}) \xleftrightarrow{P} \sigma_U^*, \quad (3a)$$

$$\sigma_D(1, 1, 1, -\frac{2}{3}, +\frac{2}{3}) \xleftrightarrow{P} \sigma_D^*, \quad (3b)$$

$$\sigma_E(1, 1, 1, -2, +2) \xleftrightarrow{P} \sigma_E^*, \quad (3c)$$

two $SU(2)$ doublets,

$$\begin{array}{c} \phi_L(1, 2, 1, -1, 0) = \begin{bmatrix} \phi_L^0 \\ \phi_L^- \end{bmatrix} \\ \uparrow P \downarrow \\ \phi_R(1, 1, 2, 0, -1) = \begin{bmatrix} \phi_R^0 \\ \phi_R^- \end{bmatrix}, \end{array} \quad (4)$$

as well as two $SU(2)$ triplets,

$$\begin{array}{c} \Delta_L(1, 3, 1, +2, 0) = \begin{bmatrix} \frac{1}{\sqrt{2}}\delta_L^+ & \delta_L^{++} \\ \delta_L^0 & -\frac{1}{\sqrt{2}}\delta_L^+ \end{bmatrix} \\ \uparrow P \downarrow \\ \Delta_R(1, 1, 3, 0, +2) = \begin{bmatrix} \frac{1}{\sqrt{2}}\delta_R^+ & \delta_R^{++} \\ \delta_R^0 & -\frac{1}{\sqrt{2}}\delta_R^+ \end{bmatrix}. \end{array} \quad (5)$$

The allowed parity-invariant Yukawa interactions then should be

$$\begin{aligned} \mathcal{L}_Y = & -y_U(\bar{q}_L\phi_L U_L^c + \bar{q}_R\phi_R U_R^c) - f_U\sigma_U\bar{U}_L U_R \\ & -y_D(\bar{q}_L\tilde{\phi}_L D_L^c + \bar{q}_R\tilde{\phi}_R D_R^c) - f_D\sigma_D\bar{D}_L D_R \\ & -y_E(\bar{l}_L\tilde{\phi}_L E_L^c + \bar{l}_R\tilde{\phi}_R E_R^c) - f_E\sigma_E\bar{E}_L E_R \\ & -\frac{1}{2}f_\Delta(\bar{l}_L^c i\tau_2 \Delta_L l_L + \bar{l}_R^c i\tau_2 \Delta_R l_R) + \text{H.c.} \quad \text{with} \\ & f_{U,D,E} = f_{U,D,E}^\dagger, \quad f_\Delta = f_\Delta^T. \end{aligned} \quad (6)$$

For simplicity, we shall not write down the full scalar potential where the parity symmetry is softly broken. Instead, we only give the terms relevant to our demonstration,

$$\begin{aligned} V \supset & \lambda(\sigma_E^*\sigma_D^3 + \text{H.c.}) + \xi(\sigma_U\sigma_D^2 + \text{H.c.}) \\ & + \eta(\sigma_E\sigma_D\sigma_U^* + \text{H.c.}) + \mu_{\phi_L}^2\phi_L^\dagger\phi_L + \mu_{\phi_R}^2\phi_R^\dagger\phi_R \\ & + \mu_{\Delta_L}^2\text{Tr}(\Delta_L^\dagger\Delta_L) + \mu_{\Delta_R}^2\text{Tr}(\Delta_R^\dagger\Delta_R) \\ & + \rho_L(\phi_L^T i\tau_2 \Delta_L \phi_L + \text{H.c.}) + \rho_R(\phi_R^T i\tau_2 \Delta_R \phi_R + \text{H.c.}) \\ & \text{with } \mu_{\phi_L}^2 \neq \mu_{\phi_R}^2, \quad \mu_{\Delta_L}^2 \neq \mu_{\Delta_R}^2, \quad \rho_L \neq \rho_R. \end{aligned} \quad (7)$$

III. SYMMETRY BREAKING

The $[SU(2)]$ -singlet Higgs scalar σ_D is responsible for spontaneously breaking the $U(1)_L \times U(1)_R$ symmetry down to a $U(1)_{B-L}$ symmetry, i.e.

$$U(1)_L \times U(1)_R \xrightarrow{\langle\sigma_D\rangle} U(1)_{B-L} \quad \text{with } \langle\sigma_D\rangle = \frac{v_D}{\sqrt{2}}. \quad (8)$$

The other $[SU(2)]$ -singlet Higgs scalars $\sigma_{U,E}$ then will acquire their vacuum expectation values (VEVs),

$$\langle\sigma_U\rangle = \frac{v_U}{\sqrt{2}} \simeq -\frac{\xi\langle\sigma_D\rangle^2}{M_{\sigma_U}^2}, \quad (9)$$

$$\langle\sigma_E\rangle = \frac{v_E}{\sqrt{2}} \simeq -\frac{\eta\langle\sigma_D\rangle\langle\sigma_U\rangle + \lambda\langle\sigma_D\rangle\langle\sigma_U^2\rangle}{M_{\sigma_E}^2}. \quad (10)$$

We can rewrite the $[SU(2)]$ -singlet Higgs scalars as

$$\sigma_D = \frac{1}{\sqrt{2}}(v_D + H_D + iA_D), \quad (11a)$$

$$\sigma_U = \frac{1}{\sqrt{2}}(v_U + H_U + iA_U), \quad (11b)$$

$$\sigma_E = \frac{1}{\sqrt{2}}(v_E + H_E + iA_E). \quad (11c)$$

The Higgs bosons $H_{D,U,E}$ will mix with each other. One of the linear combinations of the Goldstone bosons $A_{D,U,E}$ will be eaten by the gauge boson corresponding to the $U(1)_L \times U(1)_R \rightarrow U(1)_{B-L}$ symmetry breaking.

When the $[SU(2)_R]$ -doublet Higgs scalar ϕ_R develops its VEV,

$$\langle\phi_R\rangle = \begin{bmatrix} \frac{v_R}{\sqrt{2}} \\ 0 \end{bmatrix}, \quad (12)$$

the left-right symmetry will be spontaneously broken down to the electroweak symmetry, i.e.

$$SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L} \xrightarrow{\langle \phi_R \rangle} SU(3)_c \times SU(2)_L \times U(1)_Y. \quad (13)$$

Subsequently, the $[SU(2)_L]$ -doublet Higgs scalar ϕ_L will drive the spontaneous electroweak symmetry breaking, i.e.

$$SU(3)_c \times SU(2)_L \times U(1)_Y \xrightarrow{\langle \phi_L \rangle} SU(3)_c \times U(1)_{em}$$

with $\langle \phi_L \rangle = \begin{bmatrix} \frac{v_L}{\sqrt{2}} \\ 0 \end{bmatrix}$. (14)

Furthermore, the $[SU(2)]$ -triplet Higgs scalars $\Delta_{L,R}$ will acquire the induced VEVs,

$$\langle \Delta_{L,R} \rangle = \begin{bmatrix} 0 & 0 \\ \frac{u_{L,R}}{\sqrt{2}} & 0 \end{bmatrix} \quad \text{with} \quad u_{L,R} \simeq -\frac{\rho_{L,R} v_{L,R}^2}{\sqrt{2} M_{\delta_{L,R}^0}^2}. \quad (15)$$

Obviously, the VEVs $\langle \sigma_{D,U,E} \rangle$ and $\langle \phi_{L,R} \rangle$ are all real. Since the parity symmetry has been softly broken in the scalar potential (7), the left-right and electroweak symmetry breaking can occur at different scales, i.e.

$$\langle \phi_R \rangle \gg \langle \phi_L \rangle. \quad (16)$$

IV. CHARGED FERMION MASSES AND SOLUTION TO THE STRONG CP PROBLEM

After the spontaneous symmetry breaking (8-16), we can easily find the mass terms of the charged fermions,

$$\begin{aligned} \mathcal{L} \supset & - [\bar{u}_L \quad \bar{U}_R^c] \begin{bmatrix} 0 & \frac{y_U v_L}{\sqrt{2}} \\ \frac{y_U^\dagger v_R}{\sqrt{2}} & \frac{f_U^T v_E}{\sqrt{2}} \end{bmatrix} \begin{bmatrix} u_R \\ U_L^c \end{bmatrix} \\ & - [\bar{d}_L \quad \bar{D}_R^c] \begin{bmatrix} 0 & \frac{y_D v_L}{\sqrt{2}} \\ \frac{y_D^\dagger v_R}{\sqrt{2}} & \frac{f_D^T v_D}{\sqrt{2}} \end{bmatrix} \begin{bmatrix} d_R \\ D_L^c \end{bmatrix} \\ & - [\bar{e}_L \quad \bar{E}_R^c] \begin{bmatrix} 0 & \frac{y_E v_L}{\sqrt{2}} \\ \frac{y_E^\dagger v_R}{\sqrt{2}} & \frac{f_E^T v_E}{\sqrt{2}} \end{bmatrix} \begin{bmatrix} e_R \\ E_L^c \end{bmatrix} + \text{H.c.} \end{aligned} \quad (17)$$

If the off-diagonal blocks in the above mass matrices are much lighter than the diagonal blocks, we can make use of the seesaw mechanism to give the masses of the usual charged fermions from the $SU(2)$ doublets,

$$\mathcal{L} \supset -\tilde{m}_u \bar{u}_L u_R - \tilde{m}_d \bar{d}_L d_R - \tilde{m}_e \bar{e}_L e_R + \text{H.c.} \quad (18)$$

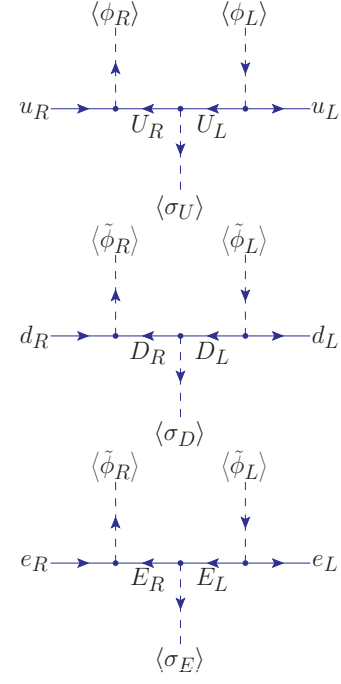


FIG. 1. Tree-level diagrams for generating the masses of up-type quarks (up), down-type quarks (middle) and charged leptons (down).

with

$$\begin{aligned} \tilde{m}_u &= -y_U \frac{v_L v_R}{2M_U} y_U^\dagger = \tilde{m}_u^\dagger \\ &= V_u \hat{m}_u V_u^\dagger = V_u \text{diag}\{m_u, m_c, m_t\} V_u^\dagger, \\ \tilde{m}_d &= -y_D \frac{v_L v_R}{2M_D} y_D^\dagger = \tilde{m}_d^\dagger \\ &= V_d \hat{m}_d V_d^\dagger = V_d \text{diag}\{m_d, m_s, m_b\} V_d^\dagger, \\ \tilde{m}_e &= -y_E \frac{v_L v_R}{2M_E} y_E^\dagger = \tilde{m}_e^\dagger \\ &= U_e \hat{m}_e U_e^\dagger = U_e \text{diag}\{m_e, m_\mu, m_\tau\} U_e^\dagger. \end{aligned} \quad (19)$$

Here, $M_{U,D,E}$ are the mass matrices of the additional charged fermions from the $SU(2)$ singlets,

$$\mathcal{L} \supset -M_U \bar{U}_L U_R - M_D \bar{D}_L D_R - M_E \bar{E}_L E_R + \text{H.c.}$$

with

$$M_{U(D,E)} = \frac{1}{\sqrt{2}} f_{U(D,E)} v_{U(D,E)} = M_{U(D,E)}^\dagger.$$

The seesaw scenario can be also understood from Fig. 1. Clearly, the CKM matrix should be $V = V_u V_d^\dagger$ as usual. In the following, we will work in the base where the mass matrices $M_{U,D,E}$ are real and diagonal.

Now the non-perturbative QCD Lagrangian should be

$$\mathcal{L}_{QCD} \supset \bar{\theta} \frac{g_3^2}{32\pi^2} G \tilde{G} \quad \text{with} \quad \bar{\theta} = \theta + \text{ArgDet}(M_u M_d), \quad (20)$$

where θ is the phase from the QCD Θ -vacuum while M_u and M_d are the mass matrices of the usual and additional down- and up-type quarks, respectively,

$$M_{d(u)} = \begin{bmatrix} 0 & \frac{y_{D(U)} v_L}{\sqrt{2}} \\ \frac{y_{D(U)}^\dagger v_R}{\sqrt{2}} & \frac{f_{D(U)}^T v_{D(U)}}{\sqrt{2}} \end{bmatrix}. \quad (21)$$

When the θ -term is removed as a result of the parity invariance, the real determinants $\text{Det}(M_d)$ and $\text{Det}(M_u)$ will lead to a zero $\text{ArgDet}(M_u M_d)$. We hence obtain a vanishing strong CP phase $\bar{\theta}$ at tree level.

V. NEUTRINO MASSES

At tree level, through their Yukawa couplings with the $[SU(2)]$ -triplet Higgs scalars, the left- and right-handed neutrinos will obtain their Majorana mass matrices

$$\mathcal{L} \supset -\frac{1}{2} m_\nu^\Pi \bar{\nu}_L \nu_L^c - \frac{1}{2} M_N \bar{N}_R N_R^c + \text{H.c.} \quad \text{with} \\ m_\nu^\Pi = f_\Delta \langle \Delta_L \rangle, \quad M_N = f_\Delta \langle \Delta_R \rangle. \quad (22)$$

The Dirac mass term between the left- and right-handed neutrinos is induced by a two-loop diagram mediated by the charged gauge bosons and fermions. One can roughly estimate

$$\mathcal{L} \supset -m_D \bar{\nu}_L N_R + \text{H.c.} \quad (23)$$

with

$$m_D \sim \frac{3g^4}{4(16\pi^2)^2} \frac{m_t m_b}{M_{W_R}^2} \tilde{m}_e. \quad (24)$$

The full masses of the left-handed neutrinos then should be a sum of the type-I and II seesaw,

$$m_\nu = m_\nu^\Pi - m_D \frac{1}{M_N} m_D^T. \quad (25)$$

It is easy to check that $m_D \lesssim 6 \text{ eV}$ for $M_{W_R} \gtrsim 1 \text{ TeV}$ and then $m_D/M_N \lesssim 6 \times 10^{-6}$, $m_D^2/M_N \lesssim 4 \times 10^{-5} \text{ eV}$ for $M_N \gtrsim 1 \text{ MeV}$. As a result, the mixing between the left- and right-handed neutrinos is negligible, while the neutrino mass matrix m_ν is dominated by the type-II seesaw contribution m_ν^Π . Accordingly the right-handed neutrino mass matrix M_N can be well described by the left-handed neutrino mass matrix m_ν up to an overall factor, i.e.

$$m_\nu = m_\nu^\Pi = U_\nu^* \hat{m}_\nu U_\nu^\dagger, \quad M_N = \frac{\langle \Delta_R \rangle}{\langle \Delta_L \rangle} m_\nu = U_\nu^* \hat{M}_N U_\nu^\dagger. \quad (26)$$

The PMNS matrix is defined by $U = U_e^\dagger U_\nu$ as usual. We can choose the base with the charged lepton mass matrix

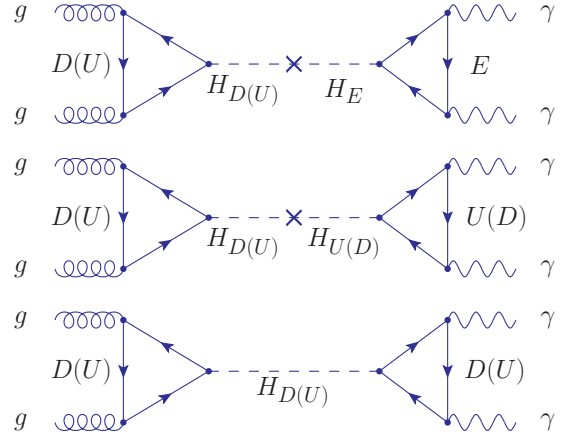


FIG. 2. The Feynman diagrams for $gg \rightarrow H \rightarrow \gamma\gamma$.

$\tilde{m}_e = \tilde{m}_e^\dagger$ being real and diagonal, and hence take $U_e = 1$ and $U_\nu = U$.

If the mixing between the left- and right-handed charged gauge bosons $W_{L,R}^\pm$ is also small, one can compare quantitatively the left-handed and right-handed neutrinoless double beta decay processes. The W_L - W_R mixing term will be induced only at one-loop level,

$$\mathcal{L} \supset \delta m_W^2 W_L^{-\mu} W_{R\mu}^+ + \text{H.c.} \quad (27)$$

with

$$\delta m_W^2 \sim \frac{3g^2}{32\pi^2} m_t m_b \sim 3 \text{ GeV}^2. \quad (28)$$

For $M_{W_R} \gtrsim 1 \text{ TeV}$, the W_L - W_R mixing will be extremely small, i.e. $\delta m_W^2/M_{W_R}^2 \lesssim 3 \times 10^{-6}$.

VI. INTERPRETATION FOR THE 750 GeV DIPHOTON RESONANCE AT ATLAS AND CMS

Recently, the ATLAS and CMS collaborations have reported a resonance of diphoton at the LHC Run-2 with energy $\sqrt{s} = 13 \text{ TeV}$ [14, 15]. The cross section $\sigma(pp \rightarrow X \rightarrow \gamma\gamma)$ is roughly estimated as $(6 - 10) \text{ fb}$, which X dubbed as the resonance. In the model the CP-even singlet Higgs $H_{D,U,E}$ mix with each other and one linear combination $H = c_U H_U + c_D H_D + c_E H_E$ could provide a good candidate for the 750 GeV resonance. For simplicity, we consider a case when H is a mixed state as

$$H = \cos \theta H_U + \sin \theta H_E. \quad (29)$$

Using the narrow width approximation, one can write the cross section $\sigma(pp \rightarrow H \rightarrow \gamma\gamma)$ through the gluon-gluon fusion channel at the 13 TeV LHC as [21]

$$\sigma(pp \rightarrow H \rightarrow \gamma\gamma) = \frac{2137}{M_H \Gamma_s} \Gamma(H \rightarrow gg) \Gamma(H \rightarrow \gamma\gamma), \quad (30)$$

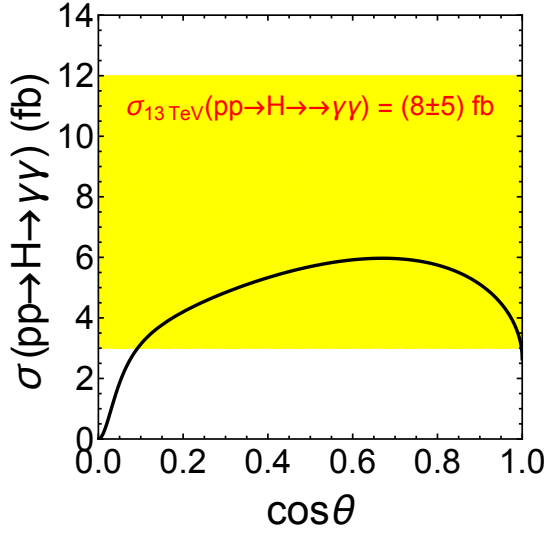


FIG. 3. The cross section $\sigma(pp \rightarrow H \rightarrow \gamma\gamma)$ as a function of the mixing $\cos\theta$. The yellow region describes the range to explain the signal observed by the ATLAS and CMS collaborations.

where M_H and Γ are the mass and width of the scalar resonance, respectively. The decay widths of $H \rightarrow \gamma\gamma$ and $H \rightarrow gg$ through the singlet fermion (U, E) loops are given by [21, 89]

$$\frac{\Gamma(H \rightarrow gg)}{M_H} = \frac{\alpha_s^2}{2\pi^3} \cos^2\theta \left(\sum_U \frac{f_U}{2\sqrt{2}} \sqrt{r_U} \mathcal{I}(r_U) \right)^2, \quad (31)$$

$$\begin{aligned} \frac{\Gamma(H \rightarrow \gamma\gamma)}{M_H} = & \frac{\alpha_{em}^2}{16\pi^3} \left[\frac{4\cos\theta}{3} \sum_U \frac{f_U}{\sqrt{2}} \sqrt{r_U} \mathcal{I}(r_U) \right. \\ & \left. + \sin\theta \sum_E \frac{f_E}{\sqrt{2}} \sqrt{r_E} \mathcal{I}(r_E) \right]^2, \end{aligned} \quad (32)$$

where $r_{U,E} = 4M_{U,E}^2/M_H^2$ and

$$\mathcal{I}(r) = 1 + (1-r) \arctan^2(1/\sqrt{r-1}). \quad (33)$$

For illustration we choose benchmark points $f_U = f_E = 2.5, v_U = 500\text{GeV}, v_E = 350\text{GeV}$, which predict the singlet fermion masses as $m_U = 884\text{ GeV}$ and $m_E = 619\text{ GeV}$. Those vector-like quark masses satisfy the current constraints imposed by the ATLAS direct search [90]. Assuming the total width is given by $\Gamma = \Gamma_{gg} + \Gamma_{\gamma\gamma}$, we plot the cross section $\sigma(pp \rightarrow H \rightarrow \gamma\gamma)$ as a function of the mixing $\cos\theta$ in Fig. 3, where $\alpha_s = 0.1$ and $\alpha_{em} = 1/128$ are used in our calculation. We also take into account of three generation vector-like quarks and leptons. The yellow region denotes the cross section needed to explain the diphoton excess. The diphoton production rate can be fulfilled at a large range of parameter space, e.g. $0.1 < \cos\theta < 0.95$. In our model the

$\Gamma(H \rightarrow Z\gamma) \sim 0.6\Gamma(H \rightarrow \gamma\gamma)$, which is well consistent with current bound [91].

VII. CONCLUSION

In summary, an $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_L \times U(1)_R \rightarrow SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ left-right symmetric framework is proposed for the universal seesaw scenario. Various vector-like $SU(2)$ -singlet fermions and Higgs scalars are introduced. The charged fermions obtain their masses through a “seesaw” way. The neutrino masses are generated through the type-II seesaw with the help of $SU(2)$ -triplet Higgs scalars, while the contribution from the canonical seesaw is negligible. The strong CP problem is solved by introducing a parity symmetry in the model. We also study the possibility of the singlet Higgs as a solution to the 750 GeV diphoton resonance observed by ATLAS and CMS at the LHC. We show that for reasonable parameters the singlet Higgs can be a good candidate for the resonance.

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